# On the 'Irreducible' Freedman-Townsend Vertex

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**Abstract**

An irreducible cohomological derivation of the Freedman-Townsend vertex in four dimensions is given.

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The problem of consistent interactions that can be introduced among gauge fields in such a way to preserve the number of gauge symmetries [1]–[4] has been reformulated as a deformation problem of the master equation [5] in the framework of the antifield BRST formalism [6]–[10]. That deformation setting was applied to Chern-Simons models [5], Yang-Mills theories [11] and two-form gauge fields [12]. The deformation procedure for two-form gauge fields employs a reducible BRST background.

The purpose of this letter is to reanalyze the problem of constructing consistent interactions among two-form gauge fields (in four dimensions), but following an *irreducible* BRST line in spite of the reducibility present within the initial model. Our method contains two basic steps. First, starting with abelian two-form gauge fields in first-order version (obtained by means of adding some auxiliary fields), we construct an irreducible BRST symmetry

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associated with the reducible one and show that we can substitute the reducible symmetry by the irreducible one. Second, we consistently deform the solution to the master equation associated with the irreducible BRST symmetry. In this manner we obtain precisely the well-known Freedman-Townsend vertex [13], the deformed gauge symmetries and also the deformed solution to the master equation. However, our deformed solution to the master equation differs from that obtained in the literature [14]–[17] by the fact that on the one hand our method introduces some additional fields (necessary at the construction of the irreducible BRST symmetry) which do not contribute to new type of couplings, and, on the other hand, the deformed BRST transformations resulting from our formalism do not involve the antifields due to the absence of terms quadratic in the antifields in the deformed solution to the master equation, so the gauge-fixed BRST symmetry is off-shell nilpotent. To our knowledge, such an irreducible procedure for the Freedman-Townsend model has not previously been published, our method establishing thus a new result.

The starting point is the Lagrangian action for abelian two-form gauge fields in first-order form (also known as the abelian Freedman-Townsend model)

$$S_0^L\left[A_\mu^a, B_a^{\mu\nu}\right] = \frac{1}{2}\int d^4x \left(-F_{\mu\nu}^a B_a^{\mu\nu} + g_{ab} A_\mu^a A^{b\mu}\right), \tag{1}$$

where $B_a^{\mu\nu}$ is an antisymmetric tensor field, the field strength of $A_\mu^a$ is defined by $F_{\mu\nu}^a = \partial_{[\mu} A_{\nu]}^a$, with $[\mu\nu]$ expressing antisymmetry with the indices between brackets, and $g_{ab}$ is an invertible, symmetric and constant matrix. It is simply to see that if we eliminate the auxiliary fields $A_\mu^a$ on their equations of motion, we recover the action of free abelian two-form gauge fields. Action (1) is invariant under the gauge transformations $\delta_\epsilon B_a^{\mu\nu} = \varepsilon^{\mu\nu\lambda\rho}\partial_\lambda \epsilon_{\rho a}$, $\delta_\epsilon A_\mu^a = 0$, where $\varepsilon^{\mu\nu\lambda\rho}$ is the antisymmetric symbol in four dimensions. The above gauge transformations are off-shell first-stage reducible as if we take $\epsilon_{\rho a} = \partial_\rho \epsilon_a$, then $\delta_\epsilon B_a^{\mu\nu} = 0$.

The reducible Lagrangian BRST symmetry corresponding to the model described by action (1), $s_R = \delta_R + \gamma_R$, contains two main pieces, namely, the Koszul-Tate differential $\delta_R$ and a model of longitudinal derivative along the gauge orbits $\gamma_R$. In the case of our model, the generators of the Koszul-Tate complex are the fermionic antighost number one antifields $B_{\mu\nu}^{*a}$ and $A_a^{*\mu}$, the bosonic antighost number two antifields $\eta^{*a\mu}$ and the fermionic antighost

number three antifields $C^{*a}$. The definitions of $\delta_R$ read as

$$\delta_R B_a^{\mu\nu} = 0,\ \delta_R A_\mu^a = 0, \tag{2}$$

$$\delta_R B_{\mu\nu}^{*a} = \frac{1}{2} F_{\mu\nu}^a,\ \delta_R A_a^{*\mu} = -\left( g_{ab} A^{b\mu} + \partial_\nu B_a^{\nu\mu} \right), \tag{3}$$

$$\delta_R \eta^{*a\mu} = \varepsilon^{\mu\nu\lambda\rho} \partial_\nu B_{\lambda\rho}^{*a}, \tag{4}$$

$$\delta_R C^{*a} = \partial_\mu \eta^{*a\mu}. \tag{5}$$

The introduction of the antifields $C^{*a}$ is implied by the necessity to 'kill' the non trivial antighost number two co-cycles $\mu^a = \partial_\mu \eta^{*a\mu}$ in the homology of $\delta_R$. The longitudinal complex contains the pure ghost number one fermionic ghosts $\eta_{a\mu}$ and the pure ghost number two bosonic ghosts for ghosts $C_a$. The definitions of $\gamma_R$ read as

$$\gamma_R A_\mu^a = 0,\ \gamma_R B_a^{\mu\nu} = \varepsilon^{\mu\nu\lambda\rho} \partial_\lambda \eta_{\rho a},\ \gamma_R \eta_{a\mu} = \partial_\mu C_a,\ \gamma_R C_a = 0. \tag{6}$$

Extending $\delta_R$ on the ghosts through $\delta_R \eta_{a\mu} = 0$, $\delta_R C_a = 0$, and $\gamma_R$ on the antifields by $\gamma_R A_a^{*\mu} = 0$, $\gamma_R B_{\mu\nu}^{*a} = 0$, $\gamma_R \eta^{*a\mu} = 0$, $\gamma_R C^{*a} = 0$, we find that $s_R^2 = 0$, $H^0 \left( s_R \right) = \{\text{physical observables}\}$, where $H^0 \left( s_R \right)$ represents the zeroth order cohomological group of $s_R$.

The main idea underlying our construction is to redefine the antifields $\eta^{*a\mu}$ in such a way that the new co-cycles of the type $\mu^a$ identically vanish. If this is done, then the antifields $C^{*a}$ are useless as there are no longer any non trivial co-cycles at antighost number two. In this way, we infer an irreducible Koszul-Tate complex, which further leads to a longitudinal complex that contains no more the ghosts for ghosts $C_a$. Accordingly our idea, we redefine the antifields $\eta^{*a\mu}$ like

$$\eta^{*a\mu} \rightarrow \eta'^{*a\mu} = M^{a\mu}_{\ \ b\nu} \eta^{*b\nu}, \tag{7}$$

where $M^{a\mu}_{\ \ b\nu}$ are taken to satisfy the conditions

$$\partial_\mu M^{a\mu}_{\ \ b\nu} = 0, \tag{8}$$

$$M^{a\mu}_{\ \ b\nu} \varepsilon^{\nu\sigma\lambda\rho} \partial_\sigma B_{\lambda\rho}^{*b} = \varepsilon^{\mu\sigma\lambda\rho} \partial_\sigma B_{\lambda\rho}^{*a}. \tag{9}$$

With the help of (4), (7) and (9) we find that

$$\delta \eta'^{*a\mu} = \varepsilon^{\mu\sigma\lambda\rho} \partial_\sigma B_{\lambda\rho}^{*a}. \tag{10}$$

The last equations do not further imply non trivial co-cycles because the new co-cycles of the type $\mu^a$ identically vanish via (8), hence we passed to an irreducible situation. In (10) we employed the notation $\delta$ instead of $\delta_R$ in order to emphasize that the Koszul-Tate complex becomes irreducible. The solution to (7–8) is expressed by

$$M^{a\mu}_{b\nu} = \delta^a{}_b \left(\delta^\mu{}_\nu - \frac{\partial^\mu \partial_\nu}{\Box}\right), \tag{11}$$

where $\Box = \partial_\lambda \partial^\lambda$. Substituting (11) in (10), we get

$$\delta \left(\eta^{*a\mu} - \frac{\partial^\mu \partial_\nu}{\Box} \eta^{*a\nu}\right) = \varepsilon^{\mu\sigma\lambda\rho} \partial_\sigma B^{*a}_{\lambda\rho}. \tag{12}$$

At this stage we introduce some scalar fields $\varphi_a$ whose antifields $\varphi^{*a}$ are demanded to be the non vanishing solutions to the equations

$$-\Box \varphi^{*a} = \delta \left(\partial_\mu \eta^{*a\mu}\right). \tag{13}$$

The non vanishing solutions $\varphi^{*a}$ enforce the irreducibility as (13) possess non vanishing solutions if and only if $\delta \left(\partial_\mu \eta^{*a\mu}\right) \neq 0$, therefore if and only if $\mu^a$ are no longer co-cycles. Using (12–13) we find that

$$\delta \eta^{*a\mu} = \varepsilon^{\mu\sigma\lambda\rho} \partial_\sigma B^{*a}_{\lambda\rho} - \partial^\mu \varphi^{*a}. \tag{14}$$

In order to preserve the nilpotency of $\delta$ we set

$$\delta \varphi^{*a} = 0. \tag{15}$$

If we maintain the actions of $\delta$ like in the reducible case

$$\delta B^{\mu\nu}_a = 0,\ \delta A^a_\mu = 0, \tag{16}$$

$$\delta B^{*a}_{\mu\nu} = \frac{1}{2} F^a_{\mu\nu},\ \delta A^{*\mu}_a = -\left(g_{ab} A^{b\mu} + \partial_\nu B^{\nu\mu}_a\right), \tag{17}$$

and define

$$\delta \varphi_a = 0, \tag{18}$$

then the formulas (14–18) describe an irreducible Koszul-Tate complex. We remark that the irreducibility was gained by introducing the supplementary

fields $\varphi_a$ and their antifields $\varphi^{*a}$ in the theory. From (14–18) we can derive the Lagrangian action and the gauge transformations of the irreducible theory. If we denote by $\tilde{S}_0^L\left[A_\mu^a, B_a^{\mu\nu}, \varphi_a\right]$ the Lagrangian action of the irreducible model, then by means of the general relations $\delta\varphi^{*a} = -\delta\tilde{S}_0^L/\delta\varphi_a$ and (15) we obtain that

$$\tilde{S}_0^L\left[A_\mu^a, B_a^{\mu\nu}, \varphi_a\right] = S_0^L\left[A_\mu^a, B_a^{\mu\nu}\right], \tag{19}$$

such that the dependence on $\varphi_a$ is trivial. On the other hand, with the help of the original gauge transformations and (14), it results that the gauge transformations of the irreducible system are expressed by

$$\delta_\epsilon B_a^{\mu\nu} = \varepsilon^{\mu\nu\lambda\rho}\partial_\lambda\epsilon_{\rho a},\ \delta_\epsilon A_\mu^a = 0,\ \delta_\epsilon\varphi_a = \partial^\mu\epsilon_{a\mu}. \tag{20}$$

In this manner, we derived an irreducible theory based on action (19) and the irreducible gauge transformations (20) associated with the abelian Freedman-Townsend model. From (19) we notice that the newly added fields $\varphi_a$ are not involved with the Lagrangian action of the irreducible theory, hence they are purely gauge. As a consequence, the physical observables (gauge invariant functions) of the irreducible model do not depend on the $\varphi_a$'s and, in addition, are invariant under the gauge transformations $\delta_\epsilon B_a^{\mu\nu} = \varepsilon^{\mu\nu\lambda\rho}\partial_\lambda\epsilon_{\rho a},\ \delta_\epsilon A_\mu^a = 0$, so they coincide with the physical observables of the original redundant theory. The construction of the irreducible longitudinal differential along the gauge orbits, $\gamma$ is realized via the definitions

$$\gamma B_a^{\mu\nu} = \varepsilon^{\mu\nu\lambda\rho}\partial_\lambda\eta_{\rho a},\ \gamma A_a^\mu = 0,\ \gamma\varphi_a = \partial^\mu\eta_{a\mu},\ \gamma\eta_{a\mu} = 0, \tag{21}$$

such that $\gamma$ is nilpotent, $\gamma^2 = 0$, without introducing the ghosts for ghosts. If we extend $\delta$ to the ghosts through $\delta\eta_{a\mu} = 0$ and $\gamma$ to the antifields by $\gamma B_{\mu\nu}^{*a} = 0,\ \gamma A_a^{*\mu} = 0,\ \gamma\varphi^{*a} = 0,\ \gamma\eta^{*a\mu} = 0$, then the homological perturbation theory [18]–[21] guarantees the existence of the irreducible BRST symmetry $s_I = \delta + \gamma$ that is nilpotent, $s_I^2 = 0$, and satisfies the property $H^0(s_I) = \{\text{physical observables}\}$, where 'physical observables' are referring to the irreducible system. As we previously mentioned, the physical observables corresponding to the reducible and irreducible formulations coincide, which leads to $H^0(s_R) = H^0(s_I)$, and moreover, the two Lagrangian BRST symmetries are nilpotent $s_R^2 = 0 = s_I^2$. By virtue of the last two relations we conclude that the two symmetries are equivalent from the BRST point of view, i.e., from the point of view of the basic equations underlying the

antifield-BRST formalism. In consequence, we can replace the reducible Lagrangian BRST symmetry with the irreducible one in the case of the model under study.

With the above conclusion at hand, we pass to the deformation procedure of the irreducible version in the context of the antifield formalism [5]. A consistent deformation of the free action $S_0^L\left[A_\mu^a, B_a^{\mu\nu}\right]$ and of its gauge invariances defines a deformation of the corresponding solution to the master equation that preserves both the master equation and the field/antifield spectra. So, if $S_0^L\left[A_\mu^a, B_a^{\mu\nu}\right]+g\int d^4x\alpha_0+O\left(g^2\right)$ stands for a consistent deformation of the free action, with deformed gauge transformations $\bar{\delta}_\epsilon B_a^{\mu\nu}=\varepsilon^{\mu\nu\lambda\rho}\partial_\lambda\epsilon_{\rho a}+g\beta_a^{\mu\nu}+O\left(g^2\right)$, $\bar{\delta}_\epsilon\varphi_a=\partial^\mu\epsilon_{a\mu}+g\beta_a+O\left(g^2\right)$, then the deformed solution to the master equation

$$\bar{S}=S+g\int d^4x\alpha+O\left(g^2\right), \tag{22}$$

satisfies $\left(\bar{S},\bar{S}\right)=0$, where

$$S=S_0^L\left[A_\mu^a, B_a^{\mu\nu}\right]+\int d^4x\left(\varepsilon^{\mu\nu\lambda\rho}B_{\mu\nu}^{*a}\partial_\lambda\eta_{\rho a}+\varphi^{*a}\partial^\mu\eta_{a\mu}\right), \tag{23}$$

and $\alpha=\alpha_0+B_{\mu\nu}^{*a}\bar{\beta}_a^{\mu\nu}+\varphi^{*a}\bar{\beta}_a+$ 'more'. Here, 'more' stands for terms 'of antighost number greater than one. The master equation $\left(\bar{S},\bar{S}\right)=0$ holds to order $g$ if and only if

$$s_I\alpha=\partial_\mu j^\mu, \tag{24}$$

for some local $j^\mu$. This means that the non trivial first-order consistent interactions belong to $H^0\left(s_I|d\right)$, where $d$ is the exterior space-time derivative. In the case where $\alpha$ is a coboundary modulo $d$ ($\alpha=s_I\rho+\partial_\mu b^\mu$), then the deformation is trivial (it can be eliminated by a redefinition of the fields). In order to investigate the solution to (24) we develop $\alpha$ accordingly the antighost number

$$\alpha=\alpha_0+\alpha_1+\ldots,\ antigh\left(\alpha_k\right)=k, \tag{25}$$

where the last term from the sum can be assumed to be annihilated by $\gamma$. Because the free theory is irreducible, we can assume that $\alpha$ stops at antighost number one, i.e., $\alpha=\alpha_0+\alpha_1$, with $\alpha_1=\alpha^{a\mu}\eta_{a\mu}$, where $\alpha^{a\mu}$ pertains to $H_1\left(\delta|d\right)$, hence is a solution of the equation $\delta\alpha^{a\mu}+\partial_\rho\lambda^{a\rho\mu}=0$. Like in the

reducible case [12], $H_2(\delta|d)$ does not vanish, but the term $\alpha_2$ can be shown to vanish. Indeed, on the one hand $\alpha_2$ is of the form $\alpha_2 = \alpha^{ab\mu\nu}\eta_{a\mu}\eta_{b\nu}$, where $\alpha^{ab\mu\nu}$ belongs to $H_2(\delta|d)$. On the other hand, the most general element in $H_2(\delta|d)$ reads as

$$\alpha^a = C^a{}_{bc}\left(\eta^{*b\mu}A^c_\mu + \frac{1}{2}\varepsilon^{\mu\nu\rho\sigma}B^{*b}_{\mu\nu}B^{*c}_{\rho\sigma} + g^{cd}\varphi^{*b}\partial_\mu A^{*\mu}_d\right), \tag{26}$$

with $g^{cd}$ the inverse of $g_{cd}$, which further gives that $\alpha^{ab\mu\nu} = \alpha^a h^{b\mu\nu}$, where $h^{b\mu\nu}$ are some constants. By Lorentz covariance $\alpha^{ab\mu\nu}$ must vanish, therefore $\alpha_2$ also vanishes. Let us investigate now the term $\alpha_1$. The general form of an object from $H_1(\delta|d)$ that is annihilated by $\gamma$ reads as

$$\alpha^{a\mu} = C^a{}_{bc}\left(\varphi^{*b}f^{c\mu}(A) + \varepsilon^{\rho\nu\lambda\mu}B^{*b}_{\rho\nu}A^c_\lambda\right), \tag{27}$$

where $f^{c\mu}(A)$ is a function of $A^a_\mu$ and $C^a{}_{bc}$ are some constants, with $C^a{}_{bc} = -C^a{}_{cb}$. It is simple to see that $\delta\alpha^{a\mu} = \partial_\rho\left(\frac{1}{2}C^a{}_{bc}\varepsilon^{\rho\nu\lambda\mu}A^b_\nu A^c_\lambda\right)$, so $\alpha^{a\mu}$ is in $H_1(\delta|d)$. On the other hand, we obtain

$$\delta\alpha_1 + \gamma\left(-\frac{1}{2}C^a{}_{bc}B^{\mu\nu}_a A^b_\mu A^c_\nu\right) = \partial_\mu\left(-\frac{1}{2}C^a{}_{bc}\varepsilon^{\mu\nu\lambda\rho}A^b_\nu A^c_\lambda\eta_{a\rho}\right). \tag{28}$$

If we compare the last equation with (24) at antighost number zero (i.e., with the equation $\delta\alpha_1 + \gamma\alpha_0 = \partial_\mu n^\mu$), it follows that

$$\alpha_0 = -\frac{1}{2}C^a{}_{bc}B^{\mu\nu}_a A^b_\mu A^c_\nu. \tag{29}$$

Thus, the deformed solution to order $g$ reads as

$$\begin{aligned}\bar{S} = S + g\int d^4x\Big(&-\frac{1}{2}C^a{}_{bc}B^{\mu\nu}_a A^b_\mu A^c_\nu + \\ &C^a{}_{bc}\left(\varphi^{*b}f^{c\mu}(A) + \varepsilon^{\rho\nu\lambda\mu}B^{*b}_{\rho\nu}A^c_\lambda\right)\eta_{a\mu}\Big).\end{aligned} \tag{30}$$

If we compute the antibracket $\left(\bar{S},\bar{S}\right)$ we obtain

$$\left(\bar{S},\bar{S}\right) = \frac{1}{3}g^2 C^e{}_{[bc}C^a{}_{d]e}\varepsilon^{\mu\nu\lambda\rho}\int d^4x A^b_\mu A^c_\nu A^d_\lambda\eta_{a\rho} \equiv g^2\int d^4x u, \tag{31}$$

where $[bcd]$ expresses the antisymmetry with respect to the indices between brackets. If we denote the term in $g^2$ from (22) by $g^2\int d^4xb$, then the interaction is consistent to order $g^2$ if and only if $u = -s_I b + \partial_\mu k^\mu$ [5]. However,

from (31) we see that $u$ cannot be of that form, and so it must vanish. This means that the constants $C^{a}{}_{bc}$ must fulfill the Jacobi identity

$$C^{e}{}_{[bc}C^{a}{}_{d]e} = 0, \tag{32}$$

hence must define the structure constants of a Lie algebra. In this situation (31) vanishes, so $\bar{S}$ (which is only of order $g$) is a solution of the master equation without adding higher order terms in $g$ (the vanishing of $u$ implies that all the higher order terms vanish).

The terms from $\bar{S}$ that do not involve the antifields, $S_0^L\left[A_\mu^a, B_a^{\mu\nu}\right] - \frac{1}{2}C^{a}{}_{bc}g\int d^4x B_a^{\mu\nu}A_\mu^b A_\nu^c$, give nothing but the well-known action of the non-abelian Freedman-Townsend model, while $\bar{S}$ itself represents the corresponding solution to the master equation deriving from our irreducible BRST approach to this model. The terms from (30) that are linear in the antifields show that the deformed gauge transformations read as $\bar{\delta}_\epsilon B_a^{\mu\nu} = \varepsilon^{\mu\nu\lambda\rho}\left(D_\lambda\right)^b{}_a \epsilon_{\rho b}$, $\bar{\delta}_\epsilon A_\mu^a = 0$, $\bar{\delta}_\epsilon \varphi_a = \partial^\mu \epsilon_{a\mu} + gC^{c}{}_{ab}f^{b\mu}\left(A\right)\epsilon_{c\mu}$, such that the gauge transformations for $B_a^{\mu\nu}$ and $A_\mu^a$ take the familiar form in the literature. In the above formulas, the covariant derivative is defined by $\left(D_\lambda\right)^b{}_a = \delta^b{}_a\partial_\lambda + gC^{b}{}_{ac}A_\lambda^c$. In addition, we have derived a class of gauge transformations for $\varphi_a$. We remark that the functions $f^{c\mu}\left(A\right)$ are still undetermined. They must be in such a way that the deformed gauge transformations are irreducible. A choice that preserves the irreducibility and, in the meantime, makes manifest the nice structure represented by the covariant derivative is $f^{c\mu}\left(A\right) = A^{c\mu}$, so $\bar{\delta}_\epsilon\varphi_a = \left(D^\mu\right)^b{}_a \epsilon_{b\mu}$. The solution (30) with $f^{c\mu}\left(A\right)$ replaced by $A^{c\mu}$ differs from that obtained in the literature by many authors [14]–[17] in the reducible framework. The solution (30) does not contain terms that are quadratic in the antifields (like in the reducible situation), so the irreducible BRST transformations $\bar{s}_I F = \left(F, \bar{S}\right)$ do not involve the antifields, such that the gauge-fixed BRST symmetry does not depend on the gauge-fixing fermion, by contrast with the reducible setting. In consequence, our approach leads to a gauge-fixed BRST symmetry that is off-shell nilpotent. Indeed, we have that $\bar{s}_I B_a^{\mu\nu} = \varepsilon^{\mu\nu\lambda\rho}\left(D_\lambda\right)^b{}_a \eta_{b\rho}$, $\bar{s}_I A_\mu^a = 0$, $\bar{s}_I\varphi_a = \left(D^\mu\right)^b{}_a \eta_{b\mu}$, $\bar{s}_I\eta_{a\mu} = 0$. In the meantime, the absence of the term quadratic in the antifields will consequently imply the absence of the three-ghost coupling term in the gauge-fixed action, such that the gauge-fixed action in the context of our irreducible approach takes a simpler form. This completes our irreducible procedure for deriving the Freedman-Townsend vertex.

To conclude with, in this letter we have exposed a cohomological approach to the Freedman-Townsend model consisting in two basic steps, namely, the construction of an irreducible BRST symmetry for the abelian version and the subsequent deformation of the irreducible theory. The results arising in our *irreducible* procedure prove the uniqueness of the Freedman-Townsend vertex in four dimensions (which has also been derived in [12], but within the reducible background) and also lead to a deformed solution of the master equation that has not previously been derived in the literature. In this light, our irreducible approach represents an efficient alternative to the reducible version exposed in [12].

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